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Remarks on T-duality for open strings†

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ABSTRACT

This contribution gives in σ-model language a short review of recent work on T-duality for open strings in the presence of abelian or non-abelian gauge fields. Furthermore, it adds a critical discussion of the relation between RG β-functions and the Born-Infeld action in the case of a string coupled to a D-brane.

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Remarks on T-duality for open strings

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This contribution gives in σ-model language a short review of recent work on T-duality for open strings in the presence of abelian or non-abelian gauge fields. Furthermore, it adds a critical discussion of the relation between RG β-functions and the Born-Infeld action in the case of a string coupled to a D-brane.

1. Introduction

Dirichlet branes, i.e. hypersurfaces which the endpoints of strings are confined to, play a fundamental role in the recent developments of the string theory duality pattern [1–3] and the connected genesis of a unique underlying theory. Such D-branes arise from open strings with free ends by T-dualizing type I string theory or have to be added to type II theories in order to fulfil all the duality requirements.

In σ-model language the T-duality transformation for open strings has been studied recently, both via formal functional integral manipulations as well as via canonical transformations [4–7].

The aim of this contribution is twofold. At first we present in section 2 a streamlined discussion of T-duality for σ-models on two-dimensional manifolds with boundary based on the functional integral formulation. It is suited for a parallel study of both abelian and non-abelian gauge fields coupling to the boundary of the manifold [5,8]. The formalism can be used to give meaning to the notion of a matrix valued D-brane position.

Secondly, we comment on a simple observation relevant for the issue of quantum equivalence of T-dual σ-models on manifolds with boundary. To give to the results gained by formal manipulations with the functional integral or by classical canonical transformations an exact status, one has to perform a careful analysis of the renormalization properties. For σ-models without boundary it has been shown that beyond 1-loop order the Buscher rules [9] require modifications [10,11]. Concerning the gauge field sector common folklore seems to take the absence of such corrections for granted. This is based on the remarkable duality property of Born-Infeld actions [12,4]. In section 3 we will point out that in the D-brane case the relation between the functional derivatives of the Born-Infeld action and the β-functions is an inhomogeneous one. This is in collision with [13] and implies the need for a more detailed study of the problem in the gauge field sector, too.

2. Functional integral formulation of T-duality

Our starting point is a generalized σ-model describing an open string moving in a target space with fields $\Psi = (G, B, \Phi)$ and A. The fields Ψ correspond to the massless excitations of closed strings. The gauge field A taking its values in the Lie algebra of some gauge group $\mathcal{G}$ corresponds to the massless excitations of open strings. To identify the dual target space fields one can restrict oneself to the partition function (A more complete study of the 2D field theory would require the inclusion of source terms for X.)

$$Z[\Psi, A] = \int DX^\mu e^{iS[\Psi;X]}\ \mathrm{tr} P e^{i\int_{\partial M} A_\mu dX^\mu}, \quad (1)$$

with

$$\begin{aligned} S[\Psi;X] &= S_M[\Psi;X] - \frac{1}{2\pi}\int_{\partial M} k(s)\Phi\ ds\ , \\ S_M[\Psi;X] &= \frac{1}{4\pi\alpha'}\int_M d^2z\sqrt{-g} \quad (2) \\ &\quad \cdot \{\partial_m X^\mu \partial_n X^\nu E^{mn}_{\mu\nu}(X(z)) \\ &\qquad + \alpha' R^{(2)}\Phi(X(z))\}, \end{aligned}$$

$$E^{mn}_{\mu\nu}(X) = g^{mn}G_{\mu\nu}(X) + \frac{\epsilon^{mn}}{\sqrt{-g}}B_{\mu\nu}(X)\ .$$

$R^{(2)}$ is the 2D curvature scalar on the manifold M, $k(s)$ the geodesic curvature of the boundary ∂M.

In the case of a non-abelian gauge group $\mathcal{G}$, due to the path ordering implied in the Wilson loop, the integrand of the functional integral appears not in the standard form as an exponential of a local action. To handle this complication we introduce an auxiliary field $\zeta(s)$, $\bar{\zeta}(s)$ living on the one-dimensional space of the variable s parametrizing the boundary ∂M of the string world sheet M. We denote its free action by S_0. It has the propagator

$$\langle \bar{\zeta}_a(s_1)\zeta_b(s_2)\rangle_0 = \delta_{ab}\Theta(s_2 - s_1) \tag{3}$$

and couples to X^μ via the interaction term

$$i\bar{\zeta}_a A_\mu^{ab}(X(z(s)))\zeta_b(s) \cdot \partial_m X^\mu(z(s)) \cdot \dot{z}^m(s) \ . \tag{4}$$

Then we can write (choosing $0 \le s \le 1$)

$$Z = \int DX^\mu \ D\bar{\zeta} \ D\zeta \ \bar{\zeta}_a(0)\zeta_a(1) \cdot \ e^{iS_0[\bar{\zeta},\zeta]} \exp(i\hat{S}[\Psi, \bar{\zeta}A\zeta; X]) \ , \tag{5}$$

with

$$\begin{aligned} \hat{S}[\Psi, C; X] \ &= \ S[\Psi; X] \\ &+ \ \int_{\partial M} C_\mu(X(z(s)), s)\dot{X}^\mu \ ds. \end{aligned} \tag{6}$$

The field

$$C_\mu(X(z), s) = \bar{\zeta}(s)A_\mu(X(z))\zeta(s) \tag{7}$$

in contrast to A_μ is no longer matrix valued. The case of an abelian gauge field is included in the formalism. However, then we can of course also drop the auxiliary field and put $C = A$. In the following we interchange the order of the X and ζ integration and adapt for our open string case the standard manipulations of the X-functional integral used to derive the T-duality rules in the closed string case [9]. Apart from the final auxiliary field integration the only difference between abelian and non-abelian A is the absence or presence of explicit s-dependence in C.

We now assume the existence of one Killing vector field $k^\mu(x)$ and invariance of our partition function under a diffeomorphism in the direction of k

$$\mathcal{L}_k\Psi \ = \ 0, \qquad \mathcal{L}_k A_\mu \ = \ D_\mu a \ . \tag{8}$$

Choosing suitable adapted coordinates we have

$$\begin{aligned} X^\mu \ &= \ (X^0, X^M), \\ k^\mu \ &= \ (1, 0), \quad \mathcal{L}_k \ = \ \partial_0 \ . \end{aligned} \tag{9}$$

Starting from an arbitrary gauge field fulfilling (8) we can arrive at $\partial_0 A_\mu = 0$ by a gauge transformation. Then all our target space fields are independent of X^0

$$\partial_0\Psi \ = \ 0, \quad \partial_0 A_\mu \ = \ 0 \ . \tag{10}$$

In the rest of the paper we take flat 2D metrics and a flat boundary and disregard the dilaton field Φ. Its transformation under T-duality should arise from some functional Jacobian as in the closed string case [9].

To perform the dualization procedure we call the result of the X-integration in (5) $\hat{Z}[\Psi, C]$ and rewrite it with the help of a Lagrange multiplier field $\tilde{X}^0$ as

$$\begin{aligned} \hat{Z}[\Psi, C] \ &= \ \int DX^\mu e^{i\hat{S}[\Psi,C;X]} \\ &= \ \int DX^M Dy_m D\tilde{X}^0 e^{i\bar{S}[\Psi,C;X^M,y_m,\tilde{X}^0]} \ , \end{aligned} \tag{11}$$

with

$$\begin{aligned} \bar{S}[\Psi, C; X^M, y_m, \tilde{X}^0] = \frac{1}{4\pi\alpha'}\int_M d^2z \\ \{\partial_m X^M \partial_n X^N E_{MN}^{mn} + 2\partial_m X^M y_n E_{M0}^{mn} \\ + y_m y_n E_{00}^{mn} + 2\tilde{X}^0 \epsilon^{mn}\partial_m y_n\} \\ + \int_{\partial M} (C_N \partial_n X^N + C_0 y_n)\dot{z}^n ds \ . \end{aligned} \tag{12}$$

The dualized version will be obtained by performing the y_m functional integral. After a shift of the integration variable suited to decouple the integral in the closed string case we get

$$\begin{aligned} \bar{S} = \frac{1}{4\pi\alpha'}\int_M d^2z\{g^{mn}G_{00}v_m v_n \\ + E_{MN}^{mn}\partial_m X^M \partial_n X^N - \tfrac{g^{mn}}{G_{00}}K_m K_n\} \\ \int_{\partial M}\{C_N \partial_n X^N \dot{z}^n + v_m L^m - \tfrac{1}{G_{00}}K_m L^m\}ds \ , \end{aligned} \tag{13}$$

with

$$K^m = E_{0N}^{mn}\partial_n X^N + \epsilon^{mn}\partial_n \tilde{X}^0 \ ,$$
$$L^m = C_0 \dot{z}^m + \frac{1}{2\pi\alpha'}\tilde{X}^0 \dot{z}^m \ ,$$
$$v_m = y_m + \frac{K_m}{G_{00}} \ . \tag{14}$$

Due to the dependence of L^m on X^M, $\tilde{X}^0$ the v integral even after the rescaling $v \to \sqrt{G_{00}}v$ does not decouple.

Since the v integral is ultralocal we can use

$$\int_M Dv_m e^{i\bar{S}} = \int_M Dv_m e^{\left(i\bar{S} - i\int_{\partial M} v_m L^m ds\right)} \cdot \int_{\partial M} Dv_m e^{\left(i\int_{\partial M} v_m L^m ds\right)} \ . \tag{15}$$

Then the v_m integral on the boundary ∂M produces a functional δ function which imposes for the remaining functional integral the constraints $L^m(z(s)) = 0, \quad m = 1, 2$, i.e.

$$2\pi\alpha' C_0(X^M(z(s)), s) \ + \ \tilde{X}^0(z(s)) = 0 \ . \tag{16}$$

Including the remaining bulk v_m integral into the overall normalization we finally arrive at (The $*$ below the integration symbol indicates the boundary condition (16).)

$$\hat{Z}[\Psi, C] \ = \ \int_* D\tilde{X}^\mu \exp(i\hat{S}[\tilde{\Psi}, \tilde{C}; \tilde{X}]) \ , \tag{17}$$

with (note $\tilde{X}^M = X^M$)

$$\tilde{X}^\mu \ = \ (\tilde{X}^0, X^M) \ , \tag{18}$$

$$\tilde{A}_\mu \ = \ (0, A_M) \ , \qquad \tilde{C}_\mu \ = \ \bar{\zeta}\tilde{A}_\mu\zeta \ , \tag{19}$$

$$\tilde{G}_{00} = \frac{1}{G_{00}} \ , \qquad \tilde{G}_{0M} = \frac{B_{0M}}{G_{00}} \ ,$$
$$\tilde{G}_{MN} = G_{MN} - \frac{G_{M0}G_{0N} + B_{M0}B_{0N}}{G_{00}} \ , \tag{20}$$

$$\tilde{B}_{0M} = \frac{G_{0M}}{G_{00}} \ ,$$
$$\tilde{B}_{MN} = B_{MN} - \frac{G_{M0}B_{0N} + B_{M0}G_{0N}}{G_{00}} \ . \tag{21}$$

For later use we denote $\hat{Z}[\Psi, C]$ understood as a functional of $\tilde{\Psi}$, $\tilde{C}$ and the function $-2\pi\alpha' C_0$ determining the boundary condition in (17) by $\mathcal{F}$

$$\hat{Z}[\Psi, C] \ = \ \mathcal{F}[\tilde{\Psi}, \tilde{C}| - 2\pi\alpha' C_0] \ . \tag{22}$$

The formulas (20, 21) coincide with those in the closed string case [9]. The gauge field $A^0(X^M)$ of the original model determines via (16) a Dirichlet boundary condition for the functional integrand in the integration defining the dual model (17). What concerns the quantized theory, the Dirichlet condition, therefore, has to be considered as an external constraint. In contrast, on the classical level there are always equivalent formulations for a boundary condition, either as an external constraint or as part of the stationarity condition of the action, possibly after a suitable addition of some boundary interaction. A discussion of T-duality and classical canonical transformations has been given in [5,7] extending corresponding results for the closed string case [14].

In the abelian case we can drop the ζ's and put $C = A$, $Z[\Psi, A] = \hat{Z}[\Psi, A]$. The boundary condition (16) has to be interpreted as a constraint confining the end points of the string to the hypersurface (Dirichlet brane)

$$\tilde{X}^0 = -2\pi\alpha' A_0(\tilde{X}^M)$$

with free movement on the brane.

In the nonabelian case (17) still has a D-brane interpretation. There is a one parameter family of D-branes

$$\tilde{X}^0 = -2\pi\alpha'\bar{\zeta}(s) A_0(\tilde{X}^M)\zeta(s) \ .$$

The endpoint of the string at string world sheet boundary parameter value s has to sit on the D-brane for parameter value s.

The final ζ-integration results in ordering the matrices sandwiched between $\bar{\zeta}(s)$ and $\zeta(s)$ with respect to increasing s. But after performing the functional integral over the world surfaces $X^\mu(z)$ there is no longer any correlation between a given target space point and s. The situation is more transparent if one treats (5) and (11) within the background field method (bm). Then both in $\mathcal{F}_{bm}$ and Z_{bm} all dependence on target space coordinates is realized via a classical string world sheet

configuration X_{cl}. The result of the ζ-integration is then [5]

$$Z_{bm}[\Psi, A; X_{cl}] = \\ \mathrm{tr} P \mathcal{F}_{bm}[\tilde{\Psi}, \tilde{A}; \tilde{X}_{cl}| - 2\pi\alpha' A_0] \ . \qquad (23)$$

Path ordering now refers to the classical path $\tilde{X}_{cl}(z(s))$ and involves both the matrices appearing in the second argument of $\mathcal{F}_{bm}$ as well as A_0 entering via the argument specifying the boundary condition.

The insertion of a matrix as boundary condition is performed *after* $\mathcal{F}_{bm}$ has been calculated with scalar (not matrix valued) s-dependent boundary condition. In this formalism we can avoid wondering about the target space interpretation of matrix valued boundaries. The situation is similar to dimensional regularization, the change from integer n to complex n is performed *after* $\int d^n x ...$ has been calculated for integer n.

3. An observation concerning quantum equivalence of T-dual models for open strings

The manipulations with the functional integral reviewed in the previous section are of pure formal nature. No attention has been paid both to functional Jacobians and renormalization issues.

In the closed string case, studied in more detail so far, there are thorough arguments for the quantum equivalence of conformal theories T-dual to one another [15]. Away from the points of conformal invariance the situation is less clear. The Buscher formulae (20), (21) can be derived by pure classical considerations on the basis of canonical transformations [14]. The Weyl anomaly coefficients, related closely to the renormalization group β-functions [16], fulfil at one-loop order the consistency equations derived from requiring (20) and (21) for the renormalized target space fields [10,17]. However, insisting on quantum equivalence in higher orders of renormalized perturbation theory requires a modification of the Buscher rules starting at two loops [10,11]. [1]

We now want to comment on some simple observations concerning quantum equivalence of the models related by (18)-(21) in the case of open strings. Then all the problems present in the absence of a boundary of the world sheet are present, too. In addition the Dirichlet boundary condition should refer at first sight to the bare quantities. For shortness we restrict ourselves to the abelian case and discuss the gauge field β-function only.

In leading order of an expansion in the derivatives of the field strength F this β-function is [18] (We absorb a factor $2\pi\alpha'$ into F. H denotes the field strength for B.)

$$\begin{aligned} \beta_\mu^{(A)} &= (B+F)_\mu{}^\nu \partial_\nu \Phi \qquad (24) \\ &+ J^{\alpha\beta}\{D_\alpha (B+F)_{\beta\mu} \\ &\qquad + \frac{1}{2}(B+F)_\mu{}^\nu H_{\nu\alpha\gamma}(B+F)^\gamma{}_\beta\} \end{aligned}$$

with

$$J^{\alpha\beta} = \left((G-(B+F)^2)^{-1}\right)^{\alpha\beta} \ . \qquad (25)$$

The vanishing of this β-function is equivalent to the stationarity condition of the generalized Born-Infeld action [18]

$$S_{\mathrm{BI}}[\Phi, G, B, F] = \int d^D x \ e^{-\Phi} \sqrt{\det(G+B+F)} \ , \qquad (26)$$

$$\frac{\delta S_{\mathrm{BI}}}{\delta A_\mu} = -e^{-\Phi}\sqrt{\det(G+B+F)} J^{\mu\nu} \cdot \beta_\nu^{(A)} \ . \qquad (27)$$

The action S_{BI} has a remarkable duality property [12,4]. For target space fields of the type discussed in the previous section ($\partial_0 \Psi = \partial_0 A = 0$) one finds by simple manipulations, based on the identity ($M_{\mu\nu}$ arbitrary matrix, $M_{00} \neq 0$)

$$\det(M_{\alpha\beta}) = M_{00} \det\left(M_{AB} - \frac{M_{A0}M_{0B}}{M_{00}}\right) ,$$

the relation

$$e^{-\Phi}\sqrt{\det(G+B+F)_{\mu\nu}} = e^{-\tilde{\Phi}}\sqrt{\det(\tilde{g}+\tilde{b}+\tilde{F})_{MN}} \ . \qquad (28)$$

[1] Repeating the formal functional integral manipulations with regularized expressions (still ignoring functional Jacobians) one expects the unmodified rules (20),(21) to hold for the bare target space fields.

Here $\tilde{g}$ and $\tilde{b}$ are the metric and the antisymmetric tensor induced on the D-brane defined by

$$f^0(X^M) = -A_0(X^M) \ , \quad f^M = X^M \ ,$$

i.e.

$$\tilde{g}_{MN} \ = \ \tilde{G}_{\mu\nu}\partial_M f^\mu \partial_N f^\nu \tag{29}$$

and similar for $\tilde{b}$, $\tilde{B}$. The fields $\tilde{G}$, $\tilde{B}$, $\tilde{A}$ are given by (19)-(21), $\tilde{\Phi} = \Phi - \frac{1}{2}\log G_{00}$.

The β-functions of a σ-model describing a string with its endpoints confined to some Dirichlet brane have been calculated in ref.[13]. Furthermore, in this paper it has been argued that the vanishing of the β-functions for the gauge field on the D-brane as well as that for the D-brane position is equivalent to the stationarity condition of the Born-Infeld action restricted to the D-brane $S_{\rm BI}[\tilde{\Phi},\tilde{g},\tilde{b},\tilde{F}]$. If this last statement was correct, due to (28) the gauge field part in the issue of quantum equivalence of T-dual models would be completely settled. However, while we agree with the result of the β-function calculation of [13] (see also [8]), the relation of these β-functions to the functional derivatives of the D-brane BI-action needs some more discussion.

The β-functions just mentioned are that for the gauge field $\tilde{A}_M$ living on the D-brane

$$\begin{aligned}\tilde{\beta}^{(\tilde{A})}_M \ &= \ (\tilde{b}+\tilde{F})_M{}^N\partial_N\tilde{\Phi} \\ &+ \ \tilde{J}^{NL}\{D_L(\tilde{b}+\tilde{F})_{NM} \\ &\qquad + \frac{1}{2}(\tilde{b}+\tilde{F})_M{}^C\tilde{h}_{CNE}(\tilde{b}+\tilde{F})^E{}_L \\ &\qquad\quad + \ (\tilde{b}+\tilde{F})_M{}^C\partial_C f^\mu K_{\mu NL}\}\end{aligned} \tag{30}$$

and that for the D-brane position f_μ

$$\begin{aligned}\tilde{\beta}^{(f)}_\mu \ &= \ -\partial_\mu\tilde{\Phi} \\ &+ \ \tilde{J}^{NL}\{\frac{1}{2}(\tilde{b}+\tilde{F})_L{}^C\tilde{H}_{\mu NC} - K_{\mu NL}\},\end{aligned} \tag{31}$$

with $\tilde{J}$ related to $\tilde{g}$, $\tilde{b}$, $\tilde{F}$ like J to G, B, F in (25) and the curvature of the D-brane given by

$$K^\mu{}_{NL} = D_N\partial_L f^\mu \ . \tag{32}$$

The two β-functions fulfil the equation

$$\begin{aligned}\tilde{\beta}^{(\tilde{A})}_M \ &= \ \tilde{J}^{NL}D_L(\tilde{b}+\tilde{F})_{NM} \\ &- \ (\tilde{b}+\tilde{F})_M{}^N\partial_N f^\mu\cdot\tilde{\beta}^{(f)}_\mu \ .\end{aligned} \tag{33}$$

For the functional derivatives of $S[\tilde{\Phi},\tilde{g},\tilde{b},\tilde{F}]$ with respect to $\tilde{A}$ and f (The last one for variations orthogonal to the D-brane.) we find

$$\begin{aligned}\frac{\delta S_{\rm BI}}{\delta\tilde{A}_M} \ &= \ -e^{\tilde{\Phi}}\sqrt{\det(\tilde{g}+\tilde{b}+\tilde{F})} \\ &\cdot \ \left(\tilde{\beta}^{(\tilde{A})}_M - \tilde{J}^{NL}(\tilde{b}+\tilde{F})_M{}^C\partial_C f^\mu K_{\mu NL}\right) \ ,\end{aligned} \tag{34}$$

$$\begin{aligned}\frac{\delta S_{\rm BI}}{\delta f^\mu_\perp} \ &= \ e^{\tilde{\Phi}}\sqrt{\det(\tilde{g}+\tilde{b}+\tilde{F})}\cdot\tilde{\beta}^{(f)}_\mu \\ &- \ 2\tilde{B}_{\mu\nu}\partial_N f^\nu\cdot\frac{\delta S_{\rm BI}}{\delta\tilde{A}_N} \ .\end{aligned} \tag{35}$$

The last two formulas disagree with ref.[13] by the second terms on their r.h.s. We want to stress that at least (34) is obvious without any calculation. The subtraction of the curvature term in (34) ensures that the gauge field derivative has the same form as in (24). This has to be correct since the gauge field $\tilde{A}_M$ is defined on the D-brane only. In contrast to $\tilde{g}$ and $\tilde{b}$ it is not induced from target space fields and does not "feel" the embedding.

Now the problem connected with (34),(35) is due to the presence of an inhomogeneity (the curvature term) in the otherwise linear relation between the β-functions and the functional derivatives of the BI-action. We no longer can claim the equivalence of the vanishing of the β-functions and the stationarity of $S_{\rm BI}$.

As far as one is interested in the equivalence of BI stationarity and the conformal invariance condition of the D-brane model there is still a way out. The Weyl anomaly coefficients are given by the β-functions plus some additive terms [16,19] usually contributing in higher orders only. While the β's are not diffeomorphism invariant the Weyl anomaly coefficients are invariant under diffeomorphisms. Hence as long as these additive terms are not calculated one should state conformal invariance as the vanishing of the β-functions up to terms which can be generated by a diffeomorphism. The disturbing term in (34) has the form of a product of $(\tilde{b}+\tilde{F})$ times a vector . Therefore, it can be generated by a combination of a diffeomorphism and a gauge transformation. Of course a final statement can be made only after a calculation of the difference between the Weyl anomaly

coefficients and the corresponding β-functions has been done.

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